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Impact of bosonic decays on the search for $\tilde{t}_1$ and $\tilde{b}_1$ squarks

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Abstract

We perform a detailed study of the decays of the lighter top and bottom squarks ($\tilde{t}_1$ and $\tilde{b}_1$) in the Minimal Supersymmetric Standard Model (MSSM). We show that the decays into Higgs or gauge bosons, i.e. $\tilde{t}_1 \to \tilde{b}_1 + (H^+$ or $W^+)$ and $\tilde{b}_1 \to \tilde{t}_1 + (H^-$ or $W^-)$, can be dominant in a wide range of the MSSM parameters due to the large Yukawa couplings and mixings of $\tilde{t}$ and $\tilde{b}$. Compared to the decays into fermions, such as $\tilde{t}_1 \to t + (\tilde{\chi}^0_i$ or $\tilde{g})$ and $\tilde{t}_1 \to b + \tilde{\chi}^+_j$, these bosonic decay modes can have significantly different decay distributions. We also show that the effect of the supersymmetric QCD running of the quark and squark parameters on the $\tilde{t}_1$ and $\tilde{b}_1$ decay branching ratios is quite dramatic. These could have an important impact on the search for $\tilde{t}_1$ and $\tilde{b}_1$ and the determination of the MSSM parameters at future colliders.

In the Minimal Supersymmetric Standard Model (MSSM) [1] supersymmetric (SUSY) partners of all the Standard Model (SM) particles are introduced. In order to solve the hierarchy, fine-tuning and naturalness problems, the SUSY particle masses have to be less than O(1 TeV). Hence the discovery of all SUSY partners and the study of their properties are essential for testing the MSSM. Future colliders, such as the Large Hadron Collider (LHC), the upgraded Tevatron, e^+e^- linear colliders, and $\mu^+\mu^-$ colliders will push the discovery reach for the SUSY particles up to the TeV mass range and allow for precise measurement of the MSSM parameters.

In this article we perform a phenomenological study concerning the search for the SUSY partners of the top (t) and bottom (b) quarks (i.e. stops ($\tilde{t}$) and sbottoms ($\tilde{b}$)). These particles may have properties which are very different from those of the squarks of the other two generations due to the large top and bottom Yukawa couplings. Production and decays of $\tilde{t}$ and $\tilde{b}$ were studied in [2, 3, 4, 5]. Stops (sbottoms) have two mass eigenstates $\tilde{t}_{1,2}$ ($\tilde{b}_{1,2}$) with $m_{\tilde{t}_1} < m_{\tilde{t}_2}$ ($m_{\tilde{b}_1} < m_{\tilde{b}_2}$). Here we focus on $\tilde{t}_1$ and $\tilde{b}_1$. Like other squarks, they can decay into fermions, i. e. a quark plus a gluino ($\tilde{g}$), neutralino ($\tilde{\chi}^0_i$) or chargino ($\tilde{\chi}^\pm_j$):

$$\begin{array}{ll} \tilde{t}_1 \to t + (\tilde{g} \text{ or } \tilde{\chi}^0_i)\,, & \tilde{b}_1 \to b + (\tilde{g} \text{ or } \tilde{\chi}^0_i)\,, \\ \tilde{t}_1 \to b + \tilde{\chi}^+_j\,, & \tilde{b}_1 \to t + \tilde{\chi}^-_j\,, \end{array} \tag{1}$$

with $i = 1, ..., 4$ and $j = 1, 2$. In addition, they can also decay into bosons [2, 3], i. e. a squark plus a Higgs or gauge boson:

$$\tilde{t}_1 \to \tilde{b}_1 + (H^+ \text{ or } W^+)\,, \qquad \tilde{b}_1 \to \tilde{t}_1 + (H^- \text{ or } W^-)\,. \tag{2}$$

The decays of Eq. (2) are possible in case the mass difference between $\tilde{t}_1$ and $\tilde{b}_1$ is sufficiently large.

In the present article we extend the analysis of [2, 3]. We point out that the bosonic decays of $\tilde{t}_1$ and $\tilde{b}_1$ of Eq. (2) can be dominant in a large region of the MSSM parameter space due to large top and bottom Yukawa couplings and large $\tilde{t}$ and $\tilde{b}$ mixing parameters. This dominance of the bosonic decays over the conventional fermionic decays of Eq. (1) could have an important impact on searches for $\tilde{t}_1$ and $\tilde{b}_1$ at future colliders. An analogous study for $\tilde{t}_2$, $\tilde{b}_2$, $\tilde{\tau}_2$, and $\tilde{\nu}_\tau$ was performed in [5, 6], where $\tilde{\tau}_2$ and $\tilde{\nu}_\tau$ are the sleptons of the third generation. In addition we show that the effect of SUSY-QCD running of the quark and squark parameters [7] on the decay branching ratios of $\tilde{t}_1$ and $\tilde{b}_1$ is quite dramatic.

First we summarize the MSSM parameters in our analysis. In the MSSM the squark sector is specified by the mass matrix in the basis $(\tilde{q}_L, \tilde{q}_R)$ with $\tilde{q} = \tilde{t}$ or $\tilde{b}$ [8, 9]

$$\mathcal{M}^2_{\tilde{q}} = \begin{pmatrix} m^2_{\tilde{q}_L} & a_q m_q \\ a_q m_q & m^2_{\tilde{q}_R} \end{pmatrix} \tag{3}$$

with

$$m^2_{\tilde{q}_L} = M^2_{\tilde{Q}} + m^2_Z \cos 2\beta\, (I_3^{q_L} - e_q \sin^2\theta_W) + m^2_q, \tag{4}$$

$$m^2_{\tilde{q}_R} = M^2_{\{\tilde{U},\tilde{D}\}} + m_Z^2 \cos 2\beta \, e_q \, \sin^2\theta_W + m_q^2, \tag{5}$$

$$a_q m_q = \begin{cases} (A_t - \mu \cot\beta) \, m_t & (\tilde{q} = \tilde{t}) \\ (A_b - \mu \tan\beta) \, m_b & (\tilde{q} = \tilde{b}) \, . \end{cases} \tag{6}$$

Here I_3^q is the third component of the weak isospin and e_q the electric charge of the quark q. $M_{\tilde{Q},\tilde{U},\tilde{D}}$ and $A_{t,b}$ are soft SUSY–breaking parameters, μ is the higgsino mass parameter, and $\tan\beta = v_2/v_1$ with v_1 (v_2) being the vacuum expectation value of the Higgs field H_1^0 (H_2^0). We treat $M_{\tilde{Q},\tilde{U},\tilde{D}}$ and $A_{t,b}$ as free parameters since the ratios $M_{\tilde{U}}/M_{\tilde{Q}}$, $M_{\tilde{D}}/M_{\tilde{Q}}$ and A_t/A_b are highly model-dependent. By diagonalizing the matrix (3) one gets the mass eigenstates $\tilde{q}_1 = \tilde{q}_L \cos\theta_{\tilde{q}} + \tilde{q}_R \sin\theta_{\tilde{q}}$, $\tilde{q}_2 = -\tilde{q}_L \sin\theta_{\tilde{q}} + \tilde{q}_R \cos\theta_{\tilde{q}}$ with the mass eigenvalues $m_{\tilde{q}_1}$, $m_{\tilde{q}_2}$ ($m_{\tilde{q}_1} < m_{\tilde{q}_2}$) and the mixing angle $\theta_{\tilde{q}}$ ($-\frac{\pi}{2} < \theta_{\tilde{q}} \le \frac{\pi}{2}$). As can be seen, sizable mixing effects can be expected in the stop sector due to the large top quark mass. Likewise, $\tilde{b}_L$–$\tilde{b}_R$ mixing may be important for large $\tan\beta$.
The properties of the charginos $\tilde{\chi}_i^\pm$ ($i = 1, 2$; $m_{\tilde{\chi}_1^\pm} < m_{\tilde{\chi}_2^\pm}$) and neutralinos $\tilde{\chi}_k^0$ ($k = 1, ..., 4$; $m_{\tilde{\chi}_1^0} < ... < m_{\tilde{\chi}_4^0}$) are determined by the parameters M, M', μ and $\tan\beta$, where M and M' are the SU(2) and U(1) gaugino mass parameters, respectively. Assuming gaugino mass unification we take $M' = (5/3)\tan^2\theta_W M$ and $m_{\tilde{g}} = (\alpha_s(m_{\tilde{g}})/\alpha_2) M$ with $m_{\tilde{g}}$ being the gluino mass. The masses and couplings of the Higgs bosons h^0, H^0, A^0 and $H^\pm$, including leading radiative corrections, are fixed by m_A, $\tan\beta$, μ, m_t, m_b, $M_{\tilde{Q}}$, $M_{\tilde{U}}$, $M_{\tilde{D}}$, A_t, and A_b. H^0 (h^0) and A^0 are the heavier (lighter) CP–even and CP–odd neutral Higgs bosons, respectively. For the radiative corrections to the h^0 and H^0 masses and their mixing angle α we use the formulae of Ref. [10]; for those to m_{H^+} we follow Ref. [11] [1].

The widths of the squark decays into Higgs and gauge bosons ($H^\pm$ and $W^\pm$) are given by [2]:

$$\Gamma(\tilde{q}_1 \to \tilde{q}'_1 \, H^\pm) = \frac{\kappa_{H^+}}{16\pi \, m^3_{\tilde{q}_1}} \, G^2_{H^+}, \tag{7}$$

$$\Gamma(\tilde{q}_1 \to \tilde{q}'_1 \, W^\pm) = \frac{\kappa^3_{W^+}}{16\pi \, m_W^2 \, m^3_{\tilde{q}_1}} \, C^2_{W^+}. \tag{8}$$

Here $(\tilde{q}_1, \tilde{q}'_1) = (\tilde{t}_1, \tilde{b}_1)$ or $(\tilde{b}_1, \tilde{t}_1)$. $\kappa_X \equiv \kappa(m^2_{\tilde{q}_1}, m^2_{\tilde{q}'_1}, m^2_X)$ is the usual kinematic factor, $\kappa(x, y, z) = (x^2 + y^2 + z^2 - 2xy - 2xz - 2yz)^{1/2}$. Notice an extra factor κ^2/m_W^2 for the gauge boson mode. The G_{H^+} denotes the squark coupling to H^+ and C_{W^+} that to W^+. Their complete expressions, as well as the widths of the fermionic modes, are given in [2, 3].

From Eqs.(3-6) we see that $m_{\tilde{t}_{1,2}} \sim M_{\tilde{Q},\tilde{U}}$ and $m_{\tilde{b}_{1,2}} \sim M_{\tilde{Q},\tilde{D}}$ in case $M_{\tilde{Q},\tilde{U},\tilde{D}}$ are large relatively to the other parameters. In this case, for $M_{\tilde{U}} > M_{\tilde{Q}} \gg M_{\tilde{D}}$ ($M_{\tilde{D}} > M_{\tilde{Q}} \gg M_{\tilde{U}}$) we have $m_{\tilde{t}_1} \gg m_{\tilde{b}_1}$ ($m_{\tilde{b}_1} \gg m_{\tilde{t}_1}$), which may allow the bosonic decays of Eq.(2). Hence in this article we consider two different patterns of the squark mass spectrum: $m_{\tilde{t}_1} \gg m_{\tilde{b}_1}$ with $(\tilde{t}_1,\tilde{b}_1) \sim (\tilde{t}_L,\tilde{b}_R)$ for $M_{\tilde{U}} \gg M_{\tilde{Q}} \gg M_{\tilde{D}}$, and $m_{\tilde{b}_1} \gg m_{\tilde{t}_1}$ with $(\tilde{t}_1,\tilde{b}_1) \sim (\tilde{t}_R,\tilde{b}_L)$ for

[1] Notice that [10, 11] have a sign convention for the parameter μ opposite to the one used here.

$M_{\tilde{D}} \gg M_{\tilde{Q}} \gg M_{\tilde{U}}$. Note here that in the former (latter) pattern the condition $M_{\tilde{U}} \gg M_{\tilde{Q}}$ ($M_{\tilde{D}} \gg M_{\tilde{Q}}$) ensures $\tilde{t}_1 \sim \tilde{t}_L$ ($\tilde{b}_1 \sim \tilde{b}_L$), which eventually enhances the bosonic decays of Eq.(2) as we will see below. Thus the bosonic decays considered here are basically the decays of $\tilde{t}_L$ into $\tilde{b}_R$ and $\tilde{b}_L$ into $\tilde{t}_R$.

The leading terms of the squark couplings to $H^\pm$ and $W^\pm$ are given by

$$\begin{aligned} G_{H^+} = G(\tilde{t}_1 \tilde{b}_1 H^\pm) &\sim h_t(\mu \sin\beta + A_t \cos\beta) \sin\theta_{\tilde{t}} \cos\theta_{\tilde{b}} \\ &\quad + h_b(\mu \cos\beta + A_b \sin\beta) \cos\theta_{\tilde{t}} \sin\theta_{\tilde{b}}, \qquad (9) \\ C_{W^+} = C(\tilde{t}_1 \tilde{b}_1 W^\pm) &\sim \frac{g}{\sqrt{2}} \cos\theta_{\tilde{t}} \cos\theta_{\tilde{b}}. \qquad (10) \end{aligned}$$

The Yukawa couplings $h_{t,b}$ are given as

$$h_t = g m_t / (\sqrt{2} m_W \sin\beta), \quad h_b = g m_b / (\sqrt{2} m_W \cos\beta). \qquad (11)$$

The Higgs bosons $H^\pm$ couple mainly to $\tilde{q}_L \tilde{q}'_R$ combinations. These couplings are proportional to the Yukawa couplings $h_{t,b}$ and the squark mixing parameters $A_{t,b}$ and μ, as can be seen in Eq.(9). Hence the widths of the squark decays into $H^\pm$ may be large for large $A_{t,b}$ and μ. Note here that the squark mixing angles $\theta_{\tilde{q}}$ themselves depend on $A_{t,b}$, μ and $\tan\beta$. In contrast, the gauge bosons $W^\pm$ couple only to $\tilde{q}_L \tilde{q}'_L$, which results in suppression of the decays into $W^\pm$. However, this suppression is largely compensated by the extra factor κ^2/m_W^2 in Eq.(8) which is very large for $m_{\tilde{q}_1} - m_{\tilde{q}'_1} \gg m_W$. (This factor stems from the contribution of the longitudinally polarized gauge boson radiation ($\tilde{q}_1 \to \tilde{q}'_1 W_L^\pm$).) Hence the widths of the squark decays into $W^\pm$ may be large for a sizable $\tilde{q}'_L - \tilde{q}'_R$ mixing term $a_{q'} m_{q'}$. On the other hand, the fermionic decays are not enhanced for large $A_{t,b}$ and μ. Therefore the branching ratios of the bosonic decays of Eq.(2) are expected to be large for large $A_{t,b}$ and μ if the gluino mode is kinematically forbidden. As for the $\tan\beta$ dependence, we expect that the branching ratio $B(\tilde{t}_1 \to \tilde{b}_1 + (H^+, W^+))$ increases with increasing $\tan\beta$ while $B(\tilde{b}_1 \to \tilde{t}_1 + (H^-, W^-))$ is rather insensitive to $\tan\beta$. The reason for this is as follows:

(i)$\tilde{t}_1$ decay;
As $(\tilde{t}_1, \tilde{b}_1) \sim (\tilde{t}_L, \tilde{b}_R)$ in this decay, the coupling $G_{H^+} = G(\tilde{t}_1 \tilde{b}_1 H^+) \sim G(\tilde{t}_L \tilde{b}_R H^+) \propto h_b \propto \tan\beta$. Hence we expect that the width $\Gamma(\tilde{t}_1 \to \tilde{b}_1 H^+)$ (and hence $B(\tilde{t}_1 \to \tilde{b}_1 H^+)$) increases with increasing $\tan\beta$. The coupling $C_{W^+} = C(\tilde{t}_1 \tilde{b}_1 W^+)$ increases with increasing $\tilde{b}_L - \tilde{b}_R$ mixing term $a_b m_b$ which increases with increasing $\tan\beta$. Hence we expect an increase of $B(\tilde{t}_1 \to \tilde{b}_1 W^+)$ as $\tan\beta$ increases.
(ii)$\tilde{b}_1$ decay;
As $(\tilde{t}_1, \tilde{b}_1) \sim (\tilde{t}_R, \tilde{b}_L)$ in this decay, the coupling $G_{H^-} = G(\tilde{t}_1 \tilde{b}_1 H^-) \sim G(\tilde{t}_R \tilde{b}_L H^-) \propto h_t$. Hence we expect that $B(\tilde{b}_1 \to \tilde{t}_1 H^-)$ is rather insensitive to $\tan\beta$ (for $\tan\beta \gtrsim 3$). The coupling $C_{W^-} = C(\tilde{t}_1 \tilde{b}_1 W^-)$ increases with increasing $\tilde{t}_L - \tilde{t}_R$ mixing term $a_t m_t$ which is fairly insensitive to $\tan\beta$ (for $\tan\beta \gtrsim 3$). Hence we expect a mild dependence of $B(\tilde{b}_1 \to \tilde{t}_1 W^-)$ on $\tan\beta$.

We now turn to the numerical analysis of the $\tilde{t}_1$ and $\tilde{b}_1$ decay branching ratios. We calculate the widths of all possibly important two-body decay modes of Eqs. (1) and (2). Three-body decays are negligible in this study. The widths of the $\tilde{t}_1$ and $\tilde{b}_1$ decays into $H^\pm$ receive large SUSY-QCD corrections for large $\tan\beta$ in the on-shell renormalization scheme [12]: the $O(\alpha_s)$ correction terms are often comparable to or even larger than the lowest order ones. Such large corrections make the perturbation calculation of the decay widths unreliable. In general this problem shows up in calculating rates of processes involving the bottom-Yukawa-coupling (h_b) for large $\tan\beta$ in the MSSM [7]. In Ref. [7] it is pointed out that this problem can be solved by carefully defining the relevant tree-level couplings in terms of appropriate running parameters and on-shell squark mixing angles $\theta_{\tilde{q}}$. Following Ref.[7], we calculate the tree-level widths of the $\tilde{t}_1$ and $\tilde{b}_1$ decays, Eqs.(1) and (2), by using the corresponding tree-level couplings defined in terms of the SUSY-QCD running parameters $m_q(Q)$ and $A_q(Q)$ (with the renormalization scale Q taken as the on-shell (pole) mass of the decaying squark), and the on-shell squark mixing angles $\theta_{\tilde{q}}$. For the kinematics, i.e. for the phase space factor $\kappa/m^3_{\tilde{q}_1}$ (for $\tilde{q}_1$ decay) the on-shell masses are used. We call the widths thus obtained as 'renormalization group (RG) improved tree-level widths'. Here we adopt the notations and conventions of [7]. Our input parameters are all on-shell ones except A_b which is a running one, i.e. they are $M_t, M_b, M_{\tilde{Q}}(\tilde{t}), M_{\tilde{U}}, M_{\tilde{D}}, A_t, A_b(Q), \mu, \tan\beta, m_A$, and M, with the renormaliztion scale Q at the on-shell (OS) mass of the decaying squark $m_{\tilde{q}_1 OS}$. $M_{t,b}$ are the on-shell (pole) masses of the t,b quarks. $M_{\tilde{Q}}(\tilde{q})$ is the on-shell $M_{\tilde{Q}}$ for the $\tilde{q}$ sector. Note that $M_{\tilde{Q}}(\tilde{t})$ is different from $M_{\tilde{Q}}(\tilde{b})$ by finite QCD corrections [7]. The procedure for obtaining all necessary on-shell and $\overline{DR}$-running parameters of quarks and squarks (i.e. on-shell (pole) $m_{\tilde{q}_{1,2}}$, on-shell $\theta_{\tilde{q}}$, and SUSY-QCD running $(m_q, A_t, M_{\tilde{Q}}, M_{\tilde{U}}, M_{\tilde{D}}, m_{\tilde{q}_{1,2}})$) is described in detail in [7]. We use this procedure. For the SM parameters we take M_t=175GeV, M_b=5GeV, m_Z=91.2GeV, $\sin^2\theta_W = 0.23, m_W = m_Z\cos\theta_W, \alpha(m_Z) = 1/129$, and $\alpha_s(m_Z) = 0.12$ [with the one-loop running $\alpha_s(Q) = 12\pi/((33-2n_f)\ln(Q^2/\Lambda^2_{n_f}))$, n_f being the number of quark flavors; only for the calculation of the SM running quark mass $m_q(Q)_{SM}$ from the two-loop renormalization group equations we use the two-loop running $\alpha_s(Q)$ as in [7]]. In order not to vary too many parameters we choose $M_{\tilde{Q}}(\tilde{t}) = \frac{3}{4}M_{\tilde{U}} = \frac{3}{2}M_{\tilde{D}}$ ($M_{\tilde{Q}}(\tilde{t})$ $= \frac{3}{2}M_{\tilde{U}} = \frac{3}{4}M_{\tilde{D}}$) for $\tilde{t}_1$ ($\tilde{b}_1$) decays, and $A_t = A_b(Q) \equiv A$ (where Q=$m_{\tilde{q}_1 OS}$ for $\tilde{q}_1$ decay) for simplicity. Moreover, we fix M=400GeV (i.e. $m_{\tilde{g}}$=1065GeV) and m_A=150GeV. Thus we have $M_{\tilde{Q}}(\tilde{t})$, A, μ and $\tan\beta$ as free parameters.

In the plots we impose the following conditions:

(i) $m_{\tilde{\chi}^\pm_1} > 100$ GeV, $m_{h^0} > 110$ GeV, $m_{\tilde{t}_1,\tilde{b}_1 OS} > m_{\tilde{\chi}^0_1}$, $m_{\tilde{\chi}^0_1} > 90$ GeV,

(ii) $\delta\rho\,(\tilde{t}\!-\!\tilde{b}) < 0.0012$ [13] using the formula of [14] (the constraint from electroweak $\delta\rho$ bound on $\tilde{t}$ and $\tilde{b}$), and

(iii) $A_t^2(Q) < 3\,(M^2_{\tilde{Q}}(Q) + M^2_{\tilde{U}}(Q) + m^2_{H_2})$, and $A_b^2(Q) < 3\,(M^2_{\tilde{Q}}(Q) + M^2_{\tilde{D}}(Q) + m^2_{H_1})$ with $m^2_{H_1} = (m^2_A + m^2_Z)\sin^2\beta - \frac{1}{2}\,m^2_Z$, $m^2_{H_2} = (m^2_A + m^2_Z)\cos^2\beta - \frac{1}{2}\,m^2_Z$, and $Q \sim M_{\tilde{Q}}$

(the approximate necessary condition for the tree-level vacuum stability [15]).

Conditions (i), along with $m_{\tilde{g}}$ = 1065 GeV, ensure that we satisfy the experimental bounds on $\tilde{\chi}_1^+$, $\tilde{\chi}_1^0$, h^0, $\tilde{t}_1$, $\tilde{b}_1$ and $\tilde{g}$ from LEP2 [16] and Tevatron [17]. Note that $m_{\tilde{\chi}_1^0} > 90$ GeV is imposed in order to evade the experimental bounds on $m_{\tilde{t}_1,\tilde{b}_1}$. Conditions (ii) and (iii) constrain the $\tilde{t}$ and $\tilde{b}$ mixings significantly. We note that the experimental data for the $b \to s\gamma$ decay give rather strong constraints [18] on the SUSY and Higgs parameters within the minimal supergravity model, especially for large $\tan\beta$. However, we do not impose this constraint since it strongly depends on the detailed properties of the squarks, including the generation-mixing.

In Fig.1 we plot in the A-μ plane the contours of the $\tilde{t}_1$ decay branching ratios of the Higgs boson mode $B(\tilde{t}_1 \to \tilde{b}_1 + H^+)$, the gauge boson mode $B(\tilde{t}_1 \to \tilde{b}_1 + W^+)$, and the total bosonic modes $B(\tilde{t}_1 \to \tilde{b}_1 + (H^+, W^+)) \equiv B(\tilde{t}_1 \to \tilde{b}_1 + H^+) + B(\tilde{t}_1 \to \tilde{b}_1 + W^+)$ at the RG-improved tree-level for $\tan\beta$=30 and $M_{\tilde{Q}}(\tilde{t})$=600GeV. We show also those of $B(\tilde{t}_1 \to \tilde{b}_1 + (H^+, W^+))$ at the naive (unimproved) tree-level, where all input parameters are bare ones, i.e. they are m_t(=175GeV), m_b(=5GeV), $M_{\tilde{Q}} = \frac{3}{4}M_{\tilde{U}} = \frac{3}{2}M_{\tilde{D}}$ (=600GeV), $A_t = A_b \equiv A$, μ, $\tan\beta$ (=30), m_A(=150GeV), and M(=400GeV) (see Eqs.(3)-(6)). We see that the $\tilde{t}_1$ decays into bosons are dominant in a large region of the A-μ plane, especially for large $|A|$ and/or $|\mu|$, as we expected. Comparing Fig.1.c with Fig.1.d we find that the effect of running of the quark and squark parameters $(m_q(Q), A_q(Q), M_{\tilde{Q},\tilde{U},\tilde{D}}(Q))$ is quite dramatic. Note here that the symmetry under $A \to -A$ and/or $\mu \to -\mu$, which holds at the naive tree-level, is strongly broken at the RG-improved tree-level. This is mainly due to the running effect steming from the gluino loops [7].

In Fig.2 we show the contours of the $\tilde{b}_1$ decay branching ratios $B(\tilde{b}_1 \to \tilde{t}_1 + H^-)$,$B(\tilde{b}_1 \to \tilde{t}_1 + W^-)$, and $B(\tilde{b}_1 \to \tilde{t}_1 + (H^-, W^-))$ at the RG-improved tree-level for $\tan\beta$=8 and 30, and $M_{\tilde{Q}}(\tilde{t})$=600GeV. The results are very similar to those for the $\tilde{t}_1$ decays. We have found again a dramatic effect of the running of the parameters.

In Fig.3 we show the individual branching ratios of the $\tilde{t}_1$ and $\tilde{b}_1$ decays as a function of $\tan\beta$ for $(A, \mu, M_{\tilde{Q}}(\tilde{t}))$=(-800, -700, 600)GeV and (800, 800, 600)GeV, respectively. One can see that the branching ratios of the $\tilde{t}_1$ decays into bosons increase with increasing $\tan\beta$ and become dominant for large $\tan\beta$ ($\gtrsim$ 20), while the $\tilde{b}_1$ decays into bosons are dominant in the entire range of $\tan\beta$ shown, as expected. As already explained, these $\tan\beta$ dependences of the $\tilde{t}_1$ and $\tilde{b}_1$ decays come from the increase of $(h_b, a_b m_b)$ with $\tan\beta$ and the mild dependence of $(h_t, a_t m_t)$ on $\tan\beta$, respectively.

In Fig.4 we show the $M_{\tilde{Q}}(\tilde{t})$ dependence of the $\tilde{t}_1$ and $\tilde{b}_1$ decay branching ratios for (A(GeV),μ(GeV),$\tan\beta$)=(-800, -700, 30) and (800, 800, 30), respectively. In these cases we have $(m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}_1^+}, m_{\tilde{g}})$=(198, 393, 1065)GeV and (198, 394, 1065)GeV, respectively. One can see that the bosonic modes dominate the $\tilde{t}_1$ and $\tilde{b}_1$ decays in a wide range of $M_{\tilde{Q}}(\tilde{t})$. For the $\tilde{b}_1$ decays we have obtained a similar result for (A(GeV), μ(GeV), $\tan\beta$)=(800, -800, 8). (Note that the decay into a gluino becomes dominant above its threshold.)

One can also see that the branching ratios of the bosonic decays decrease with increasing $M_{\tilde{Q}}(\tilde{t})$ which is roughly equal to the mass of the decaying squark $\tilde{q}_1$. This comes from the fact that in the large $M_{\tilde{Q}}(\tilde{t})(\sim m_{\tilde{q}_1})$ limit the decay widths of the bosonic and fermionic modes are proportional to $m_{\tilde{q}_1}^{-1}$ and $m_{\tilde{q}_1}$, respectively.

We find that the dominance of the bosonic modes is fairly insensitive to the choice of the values of m_A, M, and the ratio $A_b(Q)/A_t$. The decays into $H^\pm$ are kinematically suppressed for large m_A. However, the remaining gauge boson mode can still be dominant. For example, for $M_{\tilde{Q}}(\tilde{t})$=600GeV, A=μ=800GeV, and $\tan\beta$=30, we have $(m_{\tilde{t}_1 OS}, m_{\tilde{b}_1 OS}, m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}_1^+})$=(347, 602, 198, 394)GeV with $B(\tilde{b}_1 \to \tilde{t}_1\, H^-)$=(48, 37, 13, 0)% and $B(\tilde{b}_1 \to \tilde{t}_1\, W^-)$=(43, 51, 71, 82)% for m_A=(120, 200, 240, >250)GeV. We have also checked that our results do not change significantly if we take smaller values of M, as long as decays into a gluino are kinematically forbidden. As for the sensitivity of our results to the ratios $M_{\tilde{U},\tilde{D}}/M_{\tilde{Q}}(\tilde{t})$, we have found that for small $\tan\beta$ and large $|A|$ and $|\mu|$ the bosonic modes can dominate the $\tilde{b}_1$ decay even in case $M_{\tilde{Q}}(\tilde{t}) = M_{\tilde{U}} = M_{\tilde{D}}$; e.g. for $\tan\beta = 8$ and $M_{\tilde{Q}}(\tilde{t}) = M_{\tilde{U}} = M_{\tilde{D}} =$ 600GeV, we have obtained results similar to Fig.2.b and Fig.2.c. For example we have $B(\tilde{b}_1 \to \tilde{t}_1 + W^-) = 60\%$, $B(\tilde{b}_1 \to \tilde{t}_1 + (H^-, W^-)) = 90\%$ and $(m_{\tilde{t}_1 OS}, m_{\tilde{t}_2 OS}, m_{\tilde{b}_1 OS}, m_{\tilde{b}_2 OS}, m_{H^+}) = (388, 792, 570, 638, 166)$GeV for A=1200GeV, μ=-1300GeV, $\tan\beta = 8$, and $M_{\tilde{Q}}(\tilde{t}) = M_{\tilde{U}} = M_{\tilde{D}} =$ 600GeV. For these parameters the $\tilde{t}_L$-$\tilde{t}_R$ mixing term is large and the $\tilde{b}_L$-$\tilde{b}_R$ mixing term is small. This leads to a large mass-splitting of $\tilde{t}_1$-$\tilde{t}_2$ and a small one of $\tilde{b}_1$-$\tilde{b}_2$, which results in sufficient phase spaces for the bosonic decays of $\tilde{b}_1$.

In a complete analysis one would have to calculate the full SUSY-QCD one-loop corrections to the widths of the $\tilde{t}_1$ and $\tilde{b}_1$ decays. However, we expect that these corrections will not invalidate the dominance of the bosonic modes in the $\tilde{t}_1$ and $\tilde{b}_1$ decays in a significant portion of the MSSM parameter space as shown in our analysis. The calculation of the full corrections is beyond the scope of the present paper.

Now we discuss the signatures of the $\tilde{t}_1$ and $\tilde{b}_1$ decays. We compare the signals of the decays into bosons (Eq. (2)) with those of the decays into fermions (Eq. (1)). In principle, the final states of the bosonic decays can also be generated from fermionic decays. For example, the final particles of the decay chain

$$\tilde{t}_1 \to \tilde{b}_1 + (H^+ \text{ or } W^+) \to (b\tilde{\chi}_1^0) + (q\bar{q}') \tag{12}$$

are the same as those of

$$\tilde{t}_1 \to b + \tilde{\chi}_{1,2}^+ \to b + ((H^+ \text{ or } W^+) + \tilde{\chi}_1^0) \to b + (q\bar{q}'\tilde{\chi}_1^0). \tag{13}$$

Nevertheless, the decay distributions of the two processes (12) and (13) are in general different from each other due to the different intermediate states. For example, the b in the chain (12) tends to be softer than the b in (13). A similar argument holds for the quark pairs $q\bar{q}'$ in the decay chains. Moreover, the distribution of the missing energy-momentum carried by $\tilde{\chi}_1^0$ could be significantly different in (12) and (13) since it is emitted

from a different sparticle. Hence the possible dominance of the bosonic decays over the conventional fermionic decays could have an important impact on the search for the $\tilde{t}_1$ and $\tilde{b}_1$, and on the measurement of the MSSM parameters. Therefore, the effects of the bosonic decays should be included in the Monte Carlo studies of the $\tilde{t}_1$ and $\tilde{b}_1$ decays.

In conclusion, we have shown that the $\tilde{t}_1$ and $\tilde{b}_1$ decays into Higgs or gauge bosons, such as $\tilde{t}_1 \to \tilde{b}_1 + (H^+ \text{ or } W^+)$, can be dominant in a fairly wide MSSM parameter region with large mass difference between $\tilde{t}_1$ and $\tilde{b}_1$, large $|A_{t,b}|$ and/or $|\mu|$, and large $m_{\tilde{g}}$ (and large $\tan\beta$ for the $\tilde{t}_1$ decay) due to the large Yukawa couplings and mixings of $\tilde{t}$ and $\tilde{b}$. Compared to the fermionic decays, such as $\tilde{t}_1 \to b + \tilde{\chi}^+_{1,2}$, these bosonic decays can have significantly different decay distributions. We have also shown that the effect of the SUSY-QCD running of the quark and squark parameters on the $\tilde{t}_1$ and $\tilde{b}_1$ decays is quite dramatic. These could have an important impact on the searches for $\tilde{t}_1$ and $\tilde{b}_1$ and on the determination of the MSSM parameters at future colliders.

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References

[1] For reviews, see:
H. P. Nilles, Phys. Rep. 110 (1984) 1;
H.E. Haber and G.L. Kane, Phys. Rep. 117 (1985) 75;
R. Barbieri, Riv. Nuov. Cim. 11 (1988) 1.

[2] A. Bartl, W. Majerotto, and W. Porod, Z. Phys. C 64 (1994) 499; C 68 (1995) 518 (E).

[3] W. Porod, PhD Thesis, Univ. Vienna, 1997 (hep-ph/9804208).

[4] A. Bartl, H. Eberl, S. Kraml, W. Majerotto, W. Porod, and A. Sopczak, Z. Phys. C 76 (1997) 549.

[5] A. Bartl, H. Eberl, K. Hidaka, S. Kraml, T. Kon, W. Majerotto, W. Porod, and Y. Yamada, Phys. Lett. B 435 (1998) 118.

[6] A. Bartl, H. Eberl, K. Hidaka, S. Kraml, T. Kon, W. Majerotto, W. Porod, and Y. Yamada, Phys. Lett. B 460 (1999) 157.

[7] H. Eberl, K. Hidaka, S. Kraml, W. Majerotto, and Y. Yamada, Phys. Rev. D62 (2000) 055006.

[8] J. Ellis and S. Rudaz, Phys. Lett. B 128 (1983) 248.

[9] J. F. Gunion and H. E. Haber, Nucl. Phys. B 272 (1986) 1; B 402 (1993) 567 (E).

[10] J. Ellis, G. Ridolfi, and F. Zwirner, Phys. Lett. B 262 (1991) 477.

[11] A. Brignole, Phys. Lett. B 277 (1992) 313.

[12] A. Bartl, H. Eberl, K. Hidaka, S. Kraml, W. Majerotto, W. Porod, and Y. Yamada, Phys. Rev. D 59 (1999) 115007.

[13] G. Altarelli, R. Barbieri, and F. Caravaglios, Int. J. Mod. Phys. A 13 (1998) 1031.

[14] M. Drees and K. Hagiwara, Phys. Rev. D 42 (1990) 1709.

[15] J.P. Derendinger and C.A. Savoy, Nucl. Phys. B 237 (1984) 307.

[16] P. Igo-Kemenes, talk at the XXXth International Conference on High Energy Physics (ICHEP 2000), 27 July - 2 August 2000, Osaka, Japan. For transparencies see http://ichep2000.hep.sci.osaka-u.ac.jp; M. Maggi, talk at the same conference.

[17] J. Conway, talk at the same conference as [16]. For transparencies see http://ichep2000.hep.sci.osaka-u.ac.jp.

[18] H. Baer, M. Brhlik, D. Castaño, and X. Tata, Phys. Rev. D 58 (1998) 015007, and references therein.

Figure Captions

Figure 1: Contours of branching ratios of $\tilde{t}_1$ decays at the RG-improved tree-level in the A–μ plane for $\tan\beta = 30$, and $M_{\tilde{Q}}(\tilde{t}) = \frac{3}{4}M_{\tilde{U}} = \frac{3}{2}M_{\tilde{D}} = 600$GeV; (a)$B(\tilde{t}_1 \to \tilde{b}_1 + H^+)$, (b)$B(\tilde{t}_1 \to \tilde{b}_1 + W^+)$, and (c)$B(\tilde{t}_1 \to \tilde{b}_1 + (H^+, W^+))$. The regions outside of the dashed loops are excluded by the kinematics and/or the conditions (i) to (iii) given in the text. Contours of the corresponding branching ratio $B(\tilde{t}_1 \to \tilde{b}_1 + (H^+, W^+))$ at the naive tree-level are shown in Fig.d.

Figure 2: Contours of branching ratios of $\tilde{b}_1$ decays at the RG-improved tree-level in the A–μ plane for $\tan\beta$= 8 (a-c) and 30 (d-f), and $M_{\tilde{Q}}(\tilde{t}) = \frac{3}{2}M_{\tilde{U}} = \frac{3}{4}M_{\tilde{D}} = 600$GeV; (a,d)$B(\tilde{b}_1 \to \tilde{t}_1 + H^-)$, (b,e)$B(\tilde{b}_1 \to \tilde{t}_1 + W^-)$, and (c,f)$B(\tilde{b}_1 \to \tilde{t}_1 + (H^-, W^-))$. The regions outside of the dashed loops are excluded by the kinematics and/or the conditions (i) to (iii) given in the text.

Figure 3: $\tan\beta$ dependence of $\tilde{t}_1$ (a) and $\tilde{b}_1$ (b) decay branching ratios for $(A, \mu, M_{\tilde{Q}}(\tilde{t}))$ = (-800, -700, 600)GeV and (800, 800, 600)GeV, respectively. "$\tilde{b}_1 + H^+/W^+$" and "$\tilde{t}_1 + H^-/W^-$" refer to the sum of the Higgs and gauge boson modes. The gray areas are excluded by the conditions (i) to (iii) given in the text.

Figure 4: $M_{\tilde{Q}}(\tilde{t})$ dependence of $\tilde{t}_1$ (a) and $\tilde{b}_1$ (b) decay branching ratios for (A(GeV), μ(GeV), $\tan\beta$) = (-800, -700, 30) and (800, 800, 30), respectively. "$\tilde{b}_1 + H^+/W^+$" and "$\tilde{t}_1 + H^-/W^-$" refer to the sum of the Higgs and gauge boson modes. The gray areas are excluded by the conditions (i) to (iii) given in the text.

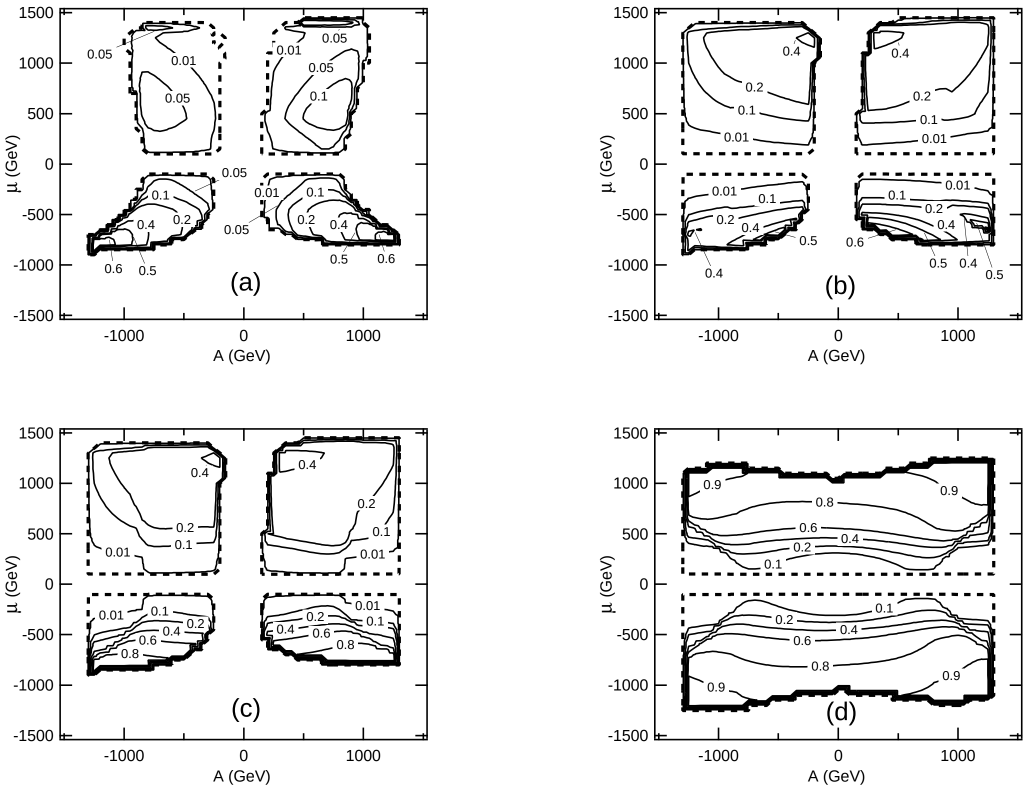


Fig.1

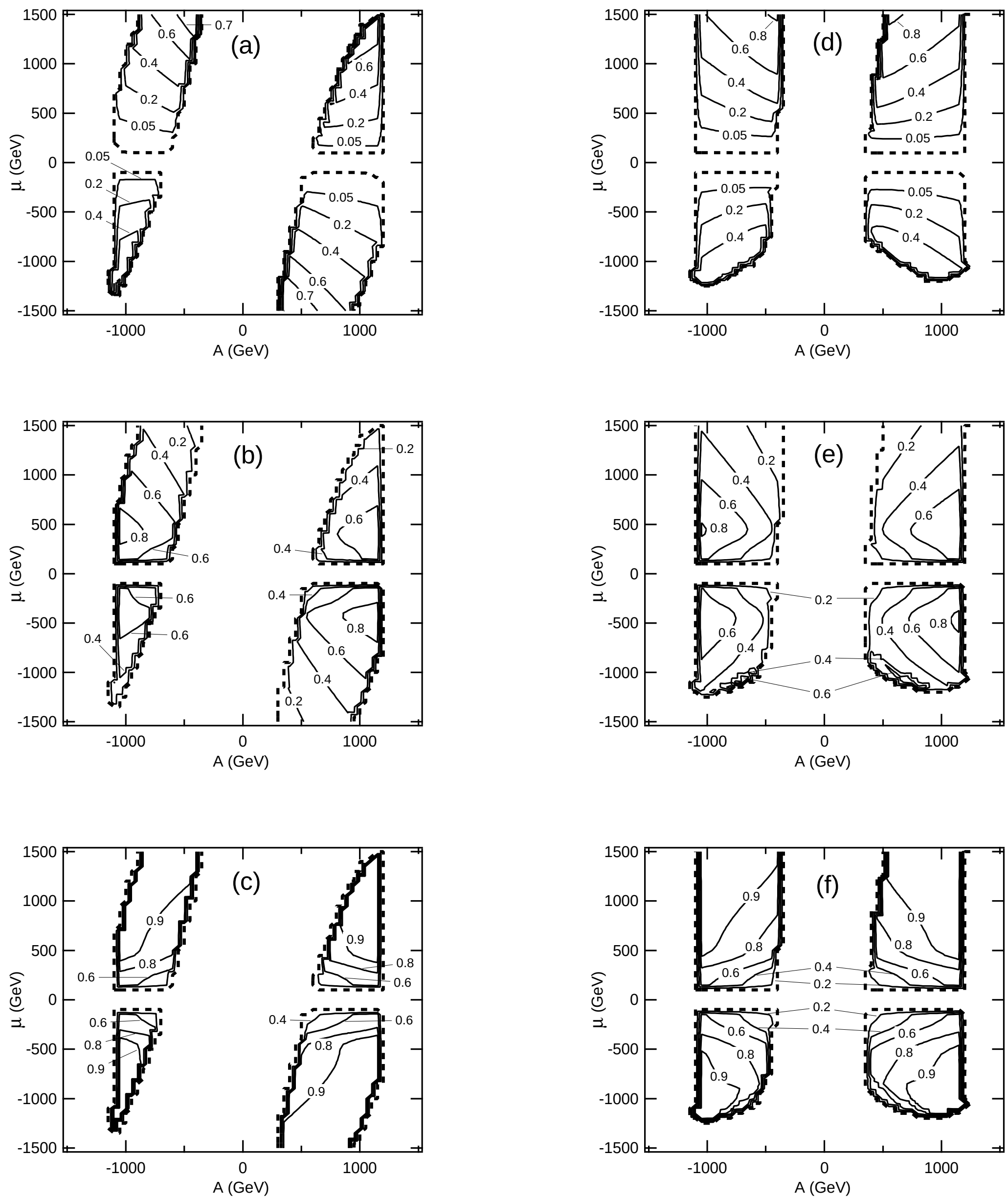
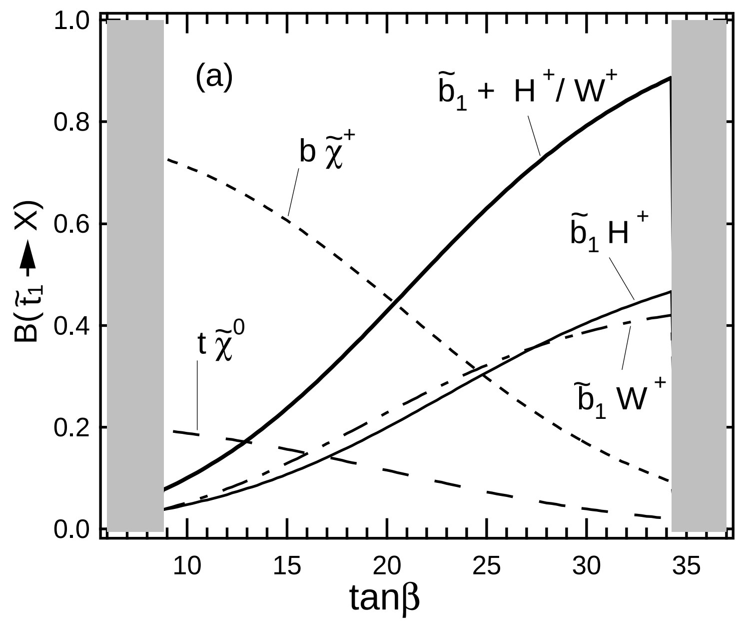
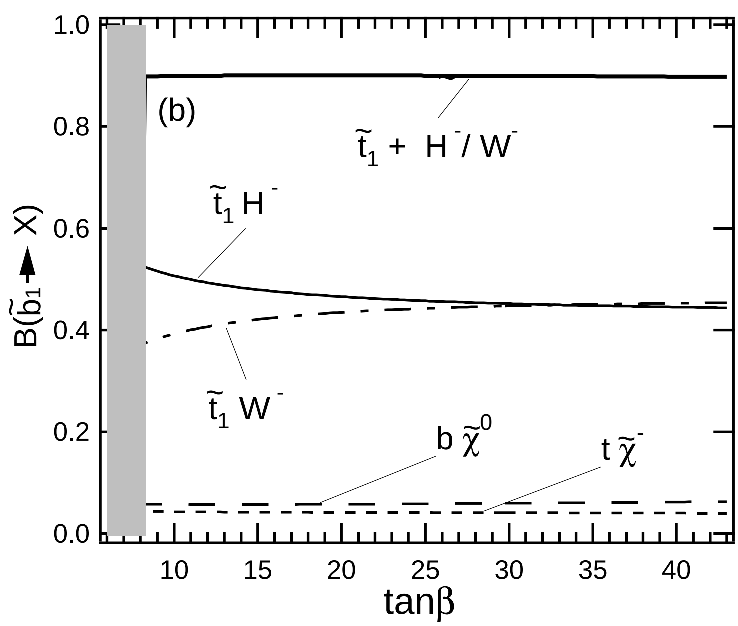


Fig.2

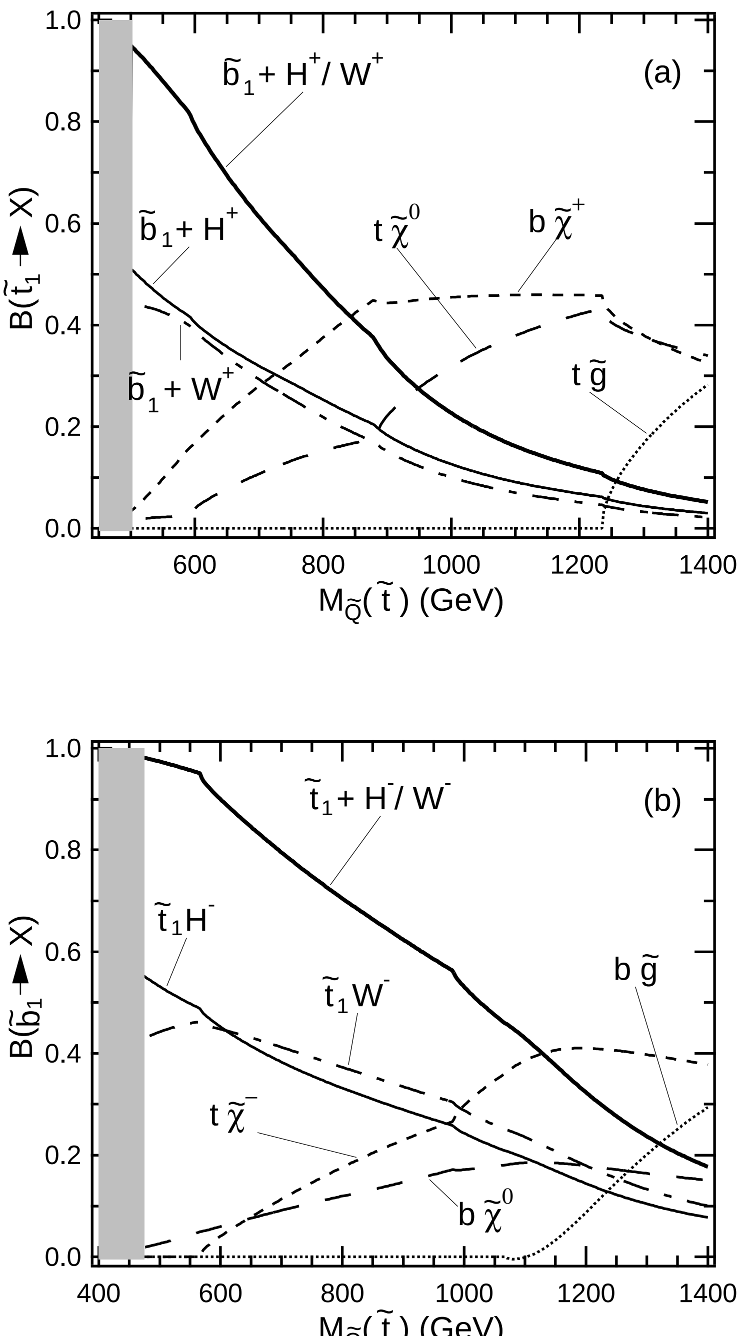


(a)
$\tilde{b}_1$ + H$^+$/ W$^+$
b $\tilde{\chi}^+$
$\tilde{b}_1$ H$^+$
t $\tilde{\chi}^0$
$\tilde{b}_1$ W$^+$
B($\tilde{t}_1$ → X)
tanβ



(b)
$\tilde{t}_1$ + H$^-$/ W$^-$
$\tilde{t}_1$ H$^-$
$\tilde{t}_1$ W$^-$
b $\tilde{\chi}^0$
t $\tilde{\chi}^-$
B($\tilde{b}_1$→ X)
tanβ

Fig.3



1.0
0.8
0.6
0.4
0.2
0.0
$B(\tilde{t}_1 \rightarrow X)$
$\tilde{b}_1 + H^+/W^+$
(a)
$\tilde{b}_1 + H^+$
$t\,\tilde{\chi}^0$
$b\,\tilde{\chi}^+$
$\tilde{b}_1 + W^+$
$t\,\tilde{g}$
600
800
1000
1200
1400
$M_{\tilde{Q}}(\tilde{t})$ (GeV)
1.0
0.8
0.6
0.4
0.2
0.0
$B(\tilde{b}_1 \rightarrow X)$
$\tilde{t}_1 + H^-/W^-$
(b)
$\tilde{t}_1 H^-$
$\tilde{t}_1 W^-$
$b\,\tilde{g}$
$t\,\tilde{\chi}^-$
$b\,\tilde{\chi}^0$
400
600
800
1000
1200
1400
$M_{\tilde{Q}}(\tilde{t})$ (GeV)

Fig.4